

# Spin Effects in Diffractive $J/\Psi$ Leptoproduction and Structure of Pomeron Coupling

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## Abstract

We calculate the cross section and the  $A_{ll}$  asymmetry of the diffractive vector meson leptoproduction for a simple model of the pomeron coupling with the proton. It is found that the sensitivity of the spin-dependent cross section of the diffractive  $J/\Psi$  production to the pomeron coupling structure is rather weak. The conclusion is made that it will be difficult to study the structure of the pomeron coupling with the proton in future polarized diffractive experiments on the  $J/\Psi$  production.

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# 1 Introduction

Study of diffractive processes at HERA has renewed interest in investigating the nature of the pomeron. New results on the pomeron intercept in diffractive events and information about the pomeron partonic structure have been obtained in H1 and ZEUS experiments [1, 2]. Among different diffractive processes, which have been studied experimentally at DESY, the vector meson production [3, 4] takes a keystone place. These reactions can give information on the gluon distribution in the nucleon at small  $x$  and on the structure of the pomeron. The diffractive  $J/\Psi$  production has a significant role in these investigations. In contrast with the  $\rho$  meson production, the  $q\bar{q}$  exchange in the  $t$ -channel is not important here and the predominant contribution is determined by a color singlet  $t$ -channel exchange (pomeron).

The phenomenological pomeron describes the cross section of elastic reactions at high energies and the low- $x$  behaviour of the structure functions. From the fit of soft elastic processes the linear soft pomeron trajectory [5] was suggested

$$\alpha_p(t) = 1 + \epsilon^s + \alpha' t, \quad (1)$$

with  $\epsilon^s = 0.08$  and  $\alpha' = 0.25\text{GeV}^{-2}$ . However, the value  $\alpha_p(0) = 1.12 - 1.2$  of the pomeron intercept, which has been obtained at HERA [2], is inconsistent with the soft pomeron with  $\epsilon^s = 0.08$ . The explanation of this discrepancy may be quite simple. In the soft reactions the interaction time is large and the pomeron rescattering effects must be important. It has been found in [6] that the rescattering effects decrease the value of the pomeron intercept from  $\epsilon \sim 0.15$  for the hard "bare" pomeron to  $\epsilon \sim 0.08$  for the soft pomeron. Thus, in hard diffractive processes like the  $J/\Psi$  production the value of  $\epsilon$  should be about 0.15.

Different models of the pomeron [7] have been used to study the diffractive vector meson production [8, 9, 10, 11]. The models [8, 9] based on the dominant role of the pomeron contribution in these processes reproduce the main features of the vector meson production: the mass,  $Q^2$ ,  $t$  and energy dependencies of the cross section. In the QCD-inspired models the pomeron is modeled by two gluons [12]. It is usually assumed that the pomeron couples to a proton like a  $C = +1$  isoscalar photon [13] and has a simple  $\gamma_\mu$  vertex. A more general form of the pomeron coupling as isoscalar nucleon current with the isoscalar Dirac and Pauli form factors have been used in [14, 15] to study the diffractive processes. In these approaches, the cross sections are dependent on the pomeron coupling with the proton. Otherwise, it has been found in [10, 11] that the cross-section of the vector-meson photoproduction in the forward limit  $|t| = 0$  and high  $Q^2$  is proportional to

$[xG(x, \bar{Q}^2)]^2$ . The typical scale here is determined by  $\bar{Q}^2 = (Q^2 + M_V^2)/4$  [10, 11] where  $Q^2$  is the photon virtuality and  $M_V$  is the mass of the vector meson. For the diffractive  $J/\Psi$  production the scale is large enough even for small  $Q^2$  and perturbative calculations can be used. We see that the cross sections of diffractive reactions are expressed, on one hand, in terms of the pomeron coupling and, on the other hand, through gluon distributions. We see, that these quantities should be related.

The sensitivity of diffractive lepto and photoproduction to the gluon density in the proton can give an excellent tool to test  $G(x)$  [10]. The relation of the spin-average diffractive production with the gluon structure function of the proton gives way to an assumption that the longitudinal double-spin asymmetry in such processes might be proportional to  $[\Delta G/G]$  [16]. Contrary to those results, in Ref. [17] it has been found that the  $A_{ll}$  asymmetry in the diffractive vector meson production should be zero for  $|t| = 0$ . As a result, this process can not be used to study  $\Delta G$  of the proton.

The value of the asymmetry in the polarized vector meson production in diffractive processes for nonzero  $|t|$  is not well known now. It is very important to perform model calculations of the spin-dependent  $J/\Psi$  leptonproduction to obtain quantitative estimations of spin asymmetries and their connection with the pomeron coupling structure (see [18]). In this paper, we calculate the cross section and the  $A_{ll}$  asymmetry of the diffractive  $J/\Psi$  leptonproduction. The cross section of the  $J/\Psi$  leptonproduction can be decomposed into the leptonic and hadronic tensors, the amplitude of the  $\gamma^* P \rightarrow J/\Psi$  transition amplitude and the pomeron exchange. After describing some kinematics of the process in Sect. 2, we consider the structure of the leptonic and hadronic tensors in Sect. 3. We use a simple form of the proton coupling with the two-gluon system which is similar to those introduced in [14]. In Sect. 4, we calculate the spin dependent cross section of the  $J/\Psi$  leptonproduction for the longitudinal polarization of the initial lepton and proton. The numerical results for the diffractive  $J/\Psi$  production at HERA and HERMES energies is presented in Sect. 5. We finish with the concluding remarks in Sect. 6.

## 2 Kinematics of Diffractive $J/\Psi$ Leptonproduction

Let us study the diffractive  $J/\Psi$  production

$$l + p \rightarrow l + p + J/\Psi \quad (2)$$

at high energies  $s = (p + l)^2$  and fixed momentum transfer  $t = r_P^2 = (p - p')^2$ . Here  $p$  and  $l$  are the initial momenta of the lepton and proton,  $p'$  is the final proton momentum and

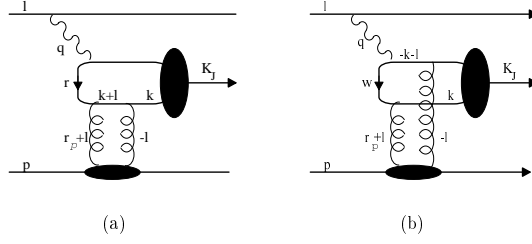


Figure 1: Feynman graphs for the diffractive vector meson production.

$r_P$  is a momentum carried by the pomeron. The graphs, which describe reaction 2, are shown in Fig. 1 a, b. The gluons from the pomeron are coupled with single and different quarks in the  $c\bar{c}$  loop. This ensures the gauge invariance of the final result [19]. The blob in the proton line represents the proton-two-gluon coupling which comprises the high energy  $t$  channel gluon ladder of the pomeron except the simple two-gluon exchange. The reaction (2) is described in addition to  $s$  and  $t$  by the variables

$$Q^2 = -q^2, \quad y = \frac{p \cdot q}{p \cdot l}, \quad x_P = \frac{q \cdot r_P}{q \cdot p}, \quad (3)$$

where  $Q^2$  is the photon virtuality.

The light-cone variables convenient for our calculations are determined by  $a_{\pm} = a_0 \pm a_z$ . In these variables the scalar production of two 4-vectors look like

$$a \cdot b = \frac{1}{2}(a_+ b_- + a_- b_+) - \vec{a}_{\perp} \vec{b}_{\perp},$$

where  $\vec{a}_{\perp}$  and  $\vec{b}_{\perp}$  represent the transverse parts of the momenta. We use the center of mass system where the momenta of the initial lepton and proton are going along the  $z$  axis and have the form

$$l = (p_+, \frac{\mu^2}{p_+}, \vec{0}), \quad p = (\frac{m^2}{p_+}, p_+, \vec{0}). \quad (4)$$

Here  $\mu$  and  $m$  are the lepton and proton mass. The energy of the lepton-proton system then reads as  $s \sim p_+^2$ . We can determine the spin vectors with positive helicity of the lepton and the proton by

$$\begin{aligned} s_l &= \frac{1}{\mu}(p_+, -\frac{\mu^2}{p_+}, \vec{0}), \quad s_l^2 = -1, \quad s_l \cdot l = 0; \\ s_p &= \frac{1}{m}(-\frac{m^2}{p_+}, p_+, \vec{0}), \quad s_p^2 = -1, \quad s_p \cdot p = 0. \end{aligned} \quad (5)$$

The momenta are carried by the photon and the pomeron and can be written as follows:

$$q = (yp_+, -\frac{Q^2}{p_+}, \vec{q}_{\perp}), \quad |q_{\perp}| = \sqrt{Q^2(1-y)}, \quad q^2 = -Q^2;$$

$$r_P = \left(-\frac{|t|}{p_+}, x_P p_+, \vec{r}_\perp\right), \quad |r_\perp| = \sqrt{|t|(1-x_P)}, \quad r_P^2 = t. \quad (6)$$

Thus,  $y$  and  $x_P$  are the fractions of the longitudinal momenta of the lepton and proton carried by the photon and pomeron, respectively. From the mass-shell equation for vector-meson momentum  $K_J^2 = (q + r_P)^2 = M_J^2$  we find that for these reactions

$$x_P \sim \frac{m_J^2 + Q^2 + |t|}{sy} \quad (7)$$

and it is small at high energies.

### 3 Structure of Leptonic and Hadronic Tensors

#### 3.1 Leptonic Tensor

The structure of the leptonic tensor is quite simple [20] because the lepton is a point-like object

$$\begin{aligned} \mathcal{L}^{\mu\nu}(s_l) &= \sum_{spin \ s_f} \bar{u}(l', s_f) \gamma^\mu u(l, s_l) \bar{u}(l, s_l) \gamma^\nu u(l', s_f) \\ &= \text{Tr} \left[ (\not{l} + \mu) \frac{1 + \gamma_5 \not{s}_l}{2} \gamma^\nu (\not{l}' + \mu) \gamma^\mu \right]. \end{aligned} \quad (8)$$

Here  $l$  and  $l'$  are the initial and final lepton momenta, and  $s_l$  is a spin vector of the initial lepton determined in (5).

The sum and difference of the cross sections with parallel and antiparallel longitudinal polarization of a proton and a lepton are expressed in terms of the spin-average and spin-dependent hadron and lepton tensors

$$\mathcal{L}^{\mu\nu}(\pm) = \frac{1}{2}(\mathcal{L}^{\mu\nu}(+\frac{1}{2}) \pm \mathcal{L}^{\mu\nu}(-\frac{1}{2})), \quad (9)$$

where  $\mathcal{L}^{\mu;\nu}(\pm\frac{1}{2})$  are the tensors with the helicity of the initial lepton equal to  $\pm 1/2$ . The tensors (9) look like

$$\begin{aligned} \mathcal{L}^{\mu\nu}(+) &= 2(g^{\mu\nu} l \cdot q + 2l^\mu l^\nu - l^\mu q^\nu - l^\nu q^\mu), \\ \mathcal{L}^{\mu\nu}(-) &= 2i\mu\epsilon^{\mu\nu\delta\rho} q_\delta(s_l)_\rho. \end{aligned} \quad (10)$$

#### 3.2 Pomeron-Proton Coupling

The pomeron is a vacuum  $t$ -channel exchange that describes diffractive processes at high energies. The hadron-hadron scattering amplitude, which is determined by the pomeron,

can be written in the form

$$T(s, t) = i\mathbb{P}(s, t)V_{A\mathbb{P}} \otimes V^{B\mathbb{P}}, \quad (11)$$

where  $\mathbb{P}$  is a function determined by the pomeron and  $V_{A\mathbb{P}}$  and  $V^{B\mathbb{P}}$  are the pomeron couplings with particles  $A$  and  $B$ , respectively.

The spin structure of the pomeron coupling is an open problem now. When the gluons from the pomeron couple to a single quark in the hadron, a simple matrix structure of the pomeron vertex

$$V_{h\mathbb{P}}^\alpha = B_{h\mathbb{P}}\gamma^\alpha \quad (12)$$

appears. This standard coupling leads to spin-flip effects decreasing with increasing energy like  $1/s$ . The effective pomeron coupling with the hadron (12) is like a  $C = +1$  isoscalar photon vertex [13]. Then, the pomeron–proton coupling should be equivalent to the isoscalar electromagnetic nucleon current. The spin-dependent pomeron coupling can be obtained if one considers together with the Dirac form factor (12) the Pauli form factor [14] in the electromagnetic nucleon current. If the gluons from the pomeron carries some fraction  $x_P$  of the initial proton momenta, this coupling can be written in the form

$$V_{pgg}^\alpha(p, t, x_P) = 2p^\alpha A(t, x_P) + \gamma^\alpha B(t, x_P). \quad (13)$$

Let us study the meson-nucleon scattering which is described at high-energies and fixed momentum transfer by the pomeron contribution. The coupling of the pomeron with the meson for small  $|t|$  can be written in the form  $q^\mu \phi(t)$  where  $q$  is the initial meson 4-momentum and  $\phi(t)$  is a meson-pomeron form factor. This form is similar to the photon–meson vertex. For the pomeron–proton coupling (13) we find the following meson-nucleon scattering amplitude

$$M(\tilde{s}, t) = i [\tilde{s}A(t, x_P) + \not{q}B(t, x_P)] \phi(t). \quad (14)$$

Here  $\tilde{s} = (p + q)^2$  and  $x_P \propto 1/\tilde{s}$ . Thus, the pomeron coupling (13) provides the standard form of the scattering amplitude. The meson–proton helicity-non-flip and helicity-flip amplitudes are expressed in terms of the invariant functions  $A$  and  $B$

$$\begin{aligned} F_{++}(s, t) &= i\tilde{s}[B(t, x_P) + 2mA(t, x_P)]\phi(t); \\ F_{+-}(s, t) &= i\tilde{s}\sqrt{|t|}A(t, x_P)\phi(t), \end{aligned} \quad (15)$$

and the spin-average cross-section is written in the form

$$\frac{d\sigma}{dt} \sim [|B + 2mA|^2 + |t||A|^2]\phi(t)^2. \quad (16)$$

Here  $m$  is a proton mass. We see that the term proportional to  $B$  represents the standard pomeron coupling that leads to the non-flip amplitude. The  $A$  function is the spin-dependent part of the pomeron coupling which produces spin-flip effects non vanishing at high-energies. The absolute value of the ratio of  $A$  to  $B$  is proportional to the ratio of helicity-flip and non-flip amplitudes. It has been found in [21, 22] that the ratio  $|A|/|B| \sim 0.1 - 0.2 \text{ GeV}^{-1}$  and has a weak energy dependence. This value of the  $|A|/|B|$  ratio and weak energy dependence of spin asymmetries in exclusive reactions is not in contradiction with the experiment [23] and is confirmed by the model results (see [21, 24] e.g.).

The proton-pomeron coupling similar to (13) has been used in [15] to analyze the spin effects in diffractive vector-meson production. In the model [21], the form (13) was found to be valid for small momentum transfer  $|t| < \text{few GeV}^2$  and the  $A(t)$  function to be caused by the meson-cloud effects in the nucleon. This model gives a quantitative description of meson-nucleon and nucleon-nucleon polarized scattering at high energies. The complicated spin structure of the pomeron coupling can be due to the nonperturbative structure of the proton. Really, in a QCD-based model, in which the proton is viewed as being composed of a quark and a diquark [25], the structure (13) for the proton coupling with a two-gluon system has been found for moderate momentum transfer [22]. The  $A(t)$  contribution is determined there by the effects of vector diquarks inside the proton, which are of an order of  $\alpha_s$ . In all the cases the spin-flip  $A(t)$  contribution is determined by the nonperturbative effects in the proton.

Since the pomeron consists of a two-gluon [12], the pomeron coupling should have two gluon indices. We use in calculations the following generalization of (13) as an ansatz:

$$\begin{aligned} V_{pgg}^{\alpha\alpha'}(p, t, x_P, l) &= 4p^\alpha p^{\alpha'} A(t, x_P, l) \\ &+ (\gamma^\alpha p^{\alpha'} + \gamma^{\alpha'} p^\alpha) B(t, x_P, l). \end{aligned} \quad (17)$$

This vertex is shown in Fig.1 a, b by the blob in the proton line. The properties of the structure (17) are completely equivalent to the coupling (13) (see Sect. 3.3 and 3.4). It has been mentioned above that this vertex contains the gluon ladder, except the two gluons which provide the  $l$  dependencies in (17). In what follows we shall calculate the imaginary part of the pomeron contribution to the scattering amplitude which dominates in the high-energy region. This contribution is equivalent to the  $t$ -channel cut in the gluon-loop graph.

### 3.3 Simple Form of Hadronic Tensor

The hadronic tensor for the vertex (13) can be written in the form

$$W^{\alpha;\beta}(s_p) = \sum_{spin\ s_f} \bar{u}(p', s_f) V_{pgg}^{\alpha}(p, t, x_P) u(p, s_p) \cdot \bar{u}(p, s_p) V_{pgg}^{\beta+}(p, t, x_P) u(p', s_f). \quad (18)$$

Here  $p$  and  $p'$  are the initial and final proton momenta and  $s_p$  is a spin vector of the initial proton. The spin-average and spin-dependent hadron tensors are determined similarly to (9). The leading term of the spin average hadron tensor must have a maximum number of the large proton momenta  $p^\alpha$ . It look like

$$W^{\alpha;\beta}(+) = 4p^\alpha p^\beta (|B + 2mA|^2 + |t||A|^2) \quad (19)$$

and is proportional to the meson-proton cross section (16) up to some function of  $t$ .

The spin-dependent part of the hadron tensor can be represented as a sum of structures which have a different nature

$$W^{\alpha;\beta}(-) = \Delta A_{full}^{\alpha;\beta} + \Delta A_1^{\alpha;\beta}. \quad (20)$$

The  $\Delta A_{full}$  term has indices of different pomeron couplings in the  $\epsilon$  function.

$$\Delta A_{full}^{\alpha;\beta} = 2im|B|^2 \epsilon^{\alpha\beta\gamma\delta} (r_P)_\gamma (s_p)_\delta \quad (21)$$

This contribution is proportional to  $|B|^2$  and equivalent in form to the spin-dependent part of the leptonic tensor (see (10)) and called by us a full block asymmetry. The  $\Delta A_1$  term contains both  $A$  and  $B$  amplitudes from (13):

$$\begin{aligned} \Delta A_1^{\alpha;\beta} &= [4p^\beta B] [2iA^* \epsilon^{\alpha\gamma\delta\rho} p_\gamma (r_P)_\delta (s_p)_\rho] \\ &- [4p^\alpha B^*] [2iA \epsilon^{\beta\gamma\delta\rho} p_\gamma (r_P)_\delta (s_p)_\rho]. \end{aligned} \quad (22)$$

### 3.4 Generalized Hadronic Tensor

The hadronic tensor for the ansatz (17) is given by

$$W^{\alpha\alpha';\beta\beta'}(s_p) = \sum_{spin\ s_f} \bar{u}(p', s_f) V_{pgg}^{\alpha\alpha'}(p, t, x_P, l) u(p, s_p) \cdot \bar{u}(p, s_p) V_{pgg}^{\beta\beta'+}(p, t, x_P, l') u(p', s_f). \quad (23)$$

The spin-average and spin-dependent hadron tensors are written as

$$W^{\alpha\alpha';\beta\beta'}(\pm) = \frac{1}{2} (W^{\alpha\alpha';\beta\beta'}(+\frac{1}{2}) \pm W^{\alpha\alpha';\beta\beta'}(-\frac{1}{2})). \quad (24)$$



For the leading term of  $W(+)$  we find

$$W^{\alpha\alpha';\beta\beta'}(+) = 16p^\alpha p^{\alpha'} p^\beta p^{\beta'} (|B + 2mA|^2 + |t||A|^2). \quad (25)$$

Note that we omit for simplicity here and in what follows the arguments of the  $A$  and  $B$  functions unless it is necessary. However, we shall remember that the amplitudes  $A$  and  $B$  depend on  $l$ , otherwise the complex conjugative values  $A^*$  and  $B^*$  are the functions of  $l'$ .

The spin-dependent part of the hadron tensor can be written as

$$W^{\alpha\alpha';\beta\beta'}(-) = \Delta A_{full}^{\alpha\alpha';\beta\beta'} + \Delta A_1^{\alpha\alpha';\beta\beta'}. \quad (26)$$

Here

$$\begin{aligned} \Delta A_{full}^{\alpha\alpha';\beta\beta'} &= 2im|B|^2 \cdot \\ &\left[ p^{\alpha'} p^{\beta'} \epsilon^{\alpha\beta\gamma\delta} (r_P)_\gamma (s_p)_\delta + \left( \text{All Permutations} \right) \left( \{ \alpha \rightarrow \alpha' \} \right) \right] \end{aligned} \quad (27)$$

and

$$\begin{aligned} \Delta A_1^{\alpha\alpha';\beta\beta'} &= \\ &\left[ 4p^\beta p^{\beta'} B \right] \left[ 2iA^* p^{\alpha'} \epsilon^{\alpha\gamma\delta\rho} p_\gamma (r_P)_\delta (s_p)_\rho + \{ \alpha \rightarrow \alpha' \} \right] - \\ &\left[ 4p^\alpha p^{\alpha'} B^* \right] \left[ 2iA p^{\beta'} \epsilon^{\beta\gamma\delta\rho} p_\gamma (r_P)_\delta (s_p)_\rho + \{ \beta \rightarrow \beta' \} \right]. \end{aligned} \quad (28)$$

The  $\Delta A_1$  term has a form of a product of the spin-dependent part  $\Delta A$  of one proton vertex to the symmetric part  $S$  of the other. Really, it can be written as

$$\Delta A_1 = -(\Delta A^{\alpha\alpha'})^* S^{\beta\beta'} - \Delta A^{\beta\beta'} (S^{\alpha\alpha'})^*, \quad (29)$$

where

$$\begin{aligned} S^{\beta\beta'} &= \left[ 4B p^\beta p^{\beta'} \right]; \\ \Delta A^{\alpha\alpha'} &= \left[ 2iA p^{\alpha'} \epsilon^{\alpha\gamma\delta\rho} p_\gamma (r_P)_\delta (s_p)_\rho + \{ \alpha \rightarrow \alpha' \} \right]. \end{aligned} \quad (30)$$

We see that  $\Delta A_{full}$  and  $\Delta A_1$  from (27) in contrast with the relevant terms in (21) have the additional  $p^{\alpha'} p^{\beta'}$  momenta and symmetrization over  $\alpha \rightarrow \alpha'$ ,  $\beta \rightarrow \beta'$  indices. The powers of large scalar production  $p \cdot q$  which appear in this case will be compensated after the loop integration over  $l$  and  $l'$ . As a result, (20) and (26) will produce the same spin-dependent amplitude. Using the mentioned argument we find that the forms (19) and (25) lead to the same spin-average amplitude too. Hence, the pomeron couplings (13) and (17) are equivalent. The very important property of (20,26) is that both the  $B^2$  and  $A \cdot B$  terms contribute to the  $W(-)$  tensor which is responsible for the asymmetry.

## 4 Diffractive $J/\Psi$ Leptoproduction

### 4.1 Amplitude of the $\gamma P \rightarrow J/\Psi$ Transition

Now we are passing to the structure of the amplitude of the  $\gamma P \rightarrow J/\Psi$  production. In what follows we have regarded the  $J/\Psi$  meson as an  $S$ -wave system of  $c\bar{c}$  quarks [26]. The  $J/\Psi$ -wave function in this case has a form  $g(\not{k} + m_c)\gamma_\mu$  where  $k$  is the momentum of quark and  $m_c$  is its mass. In the nonrelativistic approximation both the quarks have the same momenta  $k$  equal to half of the vector meson momentum  $K_J$  and the mass of  $c$  quark is equal to  $m_J/2$ . The coupling constant  $g$  can be expressed through the  $e^+e^-$  decay width of the  $J/\Psi$  meson

$$g^2 = \frac{3\Gamma_{e^+e^-}^J m_J}{64\pi\alpha^2}. \quad (31)$$

The gluons from the pomeron are coupled with the single and different quarks in the  $c\bar{c}$  loop (see Fig. 1 a, b). The  $\gamma P \rightarrow J/\Psi$  transition amplitude for these graphs look like

$$\begin{aligned} T_a &= g\text{Tr}[\not{\epsilon}_J(\not{k} + m_c)\gamma_\alpha(\not{k} + \not{l} + m_c)\gamma_{\alpha'}(\not{r} + m_c)\gamma_\nu] \frac{1}{r^2 - m_c^2}; \\ T_b &= g\text{Tr}[\not{\epsilon}_J(\not{k} + m_c)\gamma_{\alpha'}(\not{w} + m_c)\gamma_\nu(-\not{k} - \not{l} + m_c)\gamma_\alpha] \frac{1}{w^2 - m_c^2}. \end{aligned} \quad (32)$$

Here  $r = k - r_P$  and  $w = k - r_P - l$  are the momenta of the off-mass-shell quark in the loop for the diagram, Fig 1. a, b, respectively and  $e_J$  is polarization of the  $J/\Psi$  meson which obeys the relation

$$\sum_{Spin_J} e_J^\rho (e_J^\sigma)^+ = -g^{\rho\sigma} + \frac{K_J^\rho K_J^\sigma}{m_J^2}. \quad (33)$$

It is known (see [10, 11] e.g) that the leading terms of the amplitude of the diffractive vector meson production is mainly imaginary. To calculate it we must consider the  $\delta$ -function contribution in the  $s$ -channel propagators ( $k + l$  and  $p' - l$  lines for Fig 1. a). With the help of these  $\delta$  functions the integration over  $l$

$$\int d^4l = \frac{1}{2} \int dl_+ dl_- dl_\perp \quad (34)$$

can be carried out over  $l_+$  and  $l_-$  variables. One can find that both the  $l_\pm$  components of the vector  $l$  are small:  $l_+ \sim l_- \propto 1/p_+$ . This results in the transversity of the gluon momentum  $l^2 \simeq -l_\perp^2$ . The same is true for integration over  $l$  in the nonplanar graph of Fig 1. b. For the arguments in the propagator of graphs, Fig 1. a, b, we find

$$\begin{aligned} r^2 - m_c^2 &= -\frac{M_J^2 + Q^2 + |t|}{2}, \\ w^2 - m_c^2 &= -2 \left( l_\perp^2 + \vec{l}_\perp \vec{r}_\perp + \frac{M_J^2 + Q^2 + |t|}{4} \right). \end{aligned} \quad (35)$$

Thus these quark lines are far from the mass shell for heavy vector meson production even for small  $Q^2$  [10].

## 4.2 Cross Section of Vector Meson Production

The spin-average and spin dependent cross sections with parallel and antiparallel longitudinal polarization of a lepton and a proton are determined by the relation

$$\sigma(\pm) = \frac{1}{2} (\sigma(\overrightarrow{\leftarrow}) \pm \sigma(\overleftarrow{\rightarrow})). \quad (36)$$

These cross sections are expressed through the squared amplitude of the  $\gamma P \rightarrow J/\Psi$  transition convoluted with the spin-average and spin dependent lepton and hadron tensors (10), (24-26). The analyses of the leading over  $s$  contribution to the cross sections have been carried out with the help of the REDUCE and MAPLE programs. We summarize over the spin of the  $J/\Psi$  meson and use (33) in calculation. In both the cases the squared amplitude of the  $J/\Psi$  electroproduction is expressed through the integral over  $l, l'$

$$|T^\pm|^2 = \int d^2 l_\perp d^2 l'_\perp D(t, Q^2, l_\perp) D(t, Q^2, l'_\perp) \cdot F^\pm[A(l_\perp), B(l_\perp); A^*(l'_\perp), B^*(l'_\perp)], \quad (37)$$

where the functions  $F^\pm$  include the  $A$  and  $B$  amplitudes from (17). The  $l$  dependence of these functions for small  $l$  has been discussed, e.g., in the second Ref. of [9]. The function  $D$  is determined by the contribution of the  $t$ -channel gluon propagators and the sum of the  $\gamma P \rightarrow J/\Psi$  transition amplitude (32) for the graphs of Fig. 1a, b

$$D(t, Q^2, l_\perp) = \frac{1}{(l_\perp^2 + \lambda^2)((\vec{l}_\perp + \vec{r}_\perp)^2 + \lambda^2)} \cdot \left( \frac{n_a}{r^2 - m_c^2} + \frac{n_b}{w^2 - m_c^2} \right). \quad (38)$$

The leading over  $s$  terms in the numerators  $n_{a(b)}$  for the graphs, Fig. 1a, b, have a similar form but they are different in sign:  $n_b \sim -n_a = n$ . In the sum of diagrams, Fig. 1a, b, their contributions mainly compensate each other:

$$\frac{n}{w^2 - m_c^2} - \frac{n}{r^2 - m_c^2} = \frac{2n(l_\perp^2 + \vec{l}_\perp \vec{r}_\perp)}{(M_J^2 + Q^2 + |t|) [\vec{l}_\perp^2 + \vec{l}_\perp \vec{r}_\perp + (M_J^2 + Q^2 + |t|)/4]}. \quad (39)$$

This function determines the  $Q^2$ -dependence of  $D$ . It can be seen that the typical scale in the integral (37) is defined by (39) and of the order  $\vec{l}_\perp^2 \sim (M_J^2 + Q^2 + |t|)/4$  [10, 11].

As a result, (37) can be estimated in the form

$$\begin{aligned} |T^\pm|^2 &= F^\pm(A, B; A^*, B^*) I^2; \\ I &= \int d^2 l_\perp D(t, Q^2, l_\perp). \end{aligned} \quad (40)$$

The functions  $A$  and  $B$  in  $F^\pm$  are dependent on the scale

$$\bar{l}_\perp^2 \sim \bar{Q}^2 = (M_J^2 + Q^2 + |t|)/4. \quad (41)$$

The cross section of the  $J/\Psi$  leptonproduction can be written in the form

$$\frac{d\sigma^\pm}{dQ^2 dy dt} = \frac{|T^\pm|^2}{32(2\pi)^3 Q^2 s^2 y}. \quad (42)$$

For the spin-average squared amplitude we find

$$\begin{aligned} |T^+|^2 &= N((2 - 2y + y^2)m_J^2 + 2(1 - y)Q^2)s^2 \cdot \\ &[|B + 2mA|^2 + |A|^2|t|]I^2. \end{aligned} \quad (43)$$

Here  $N$  is a normalization factor

$$N = \frac{\Gamma_{e^+e^-}^J M_J \alpha_s^4}{27\pi^2}. \quad (44)$$

In (43) the term proportional to  $(2 - 2y + y^2)m_J^2$  represents the contribution of a virtual photon with transverse polarization. The  $2(1 - y)Q^2$  term describes the effect of longitudinal photons. Thus, we see that the ratio of  $\sigma_T/\sigma_L \propto m_J^2/Q^2$ . Such a behaviour is typical of a simple form of the vector meson wave function used here (see e.g. [27])

The spin-dependent squared amplitude look like

$$|T^-|^2 = N(2 - y)s|t| [|B|^2 + m(A^*B + AB^*)]m_J^2 I^2. \quad (45)$$

We see that  $|T^-|^2$  vanishes in the forward direction ( $t = 0$ ) and is suppressed as a power of  $s$  with respect to (43). The reason for this suppression is quite simple. The leading contribution to  $\sigma(-)$  is going from the term  $\epsilon^{\alpha\beta\gamma\rho}(r_P)_\gamma$ .. which is proportional to  $x_P p$ . As a result, an additional  $x_P$  appears in  $\sigma(-)$ . It has been confirmed by the calculation of the  $A_U$  asymmetry in different diffractive reactions [18, 28]. In the case of vector meson production,  $x_P$  is small (7) and behaves like  $1/s$ . Hence, longitudinal double-spin asymmetry in this diffractive process will be small at high energies which is confirmed by our calculation for (45) and (43).

## 5 Numerical Results for Spin Dependent Cross Section

We shall calculate the polarized cross section of the diffractive  $J/\Psi$  production (42) for the amplitudes (43, 45). The connection of the spin-average cross section of the  $J/\Psi$  production with the gluon distribution function is known [10, 11]

$$\left. \frac{d\sigma}{dt} \right|_{t \sim 0} \propto F_B^2(t) \left( x_P G(x_P, \bar{Q}^2) \right)^2. \quad (46)$$

Here  $F_B(t)$  is a new form factor which describes  $t$ -dependence of the two-gluon coupling with the proton. The expression of this cross section through the pomeron-proton structure has been found in (43). It can be seen that the  $B$  function in (17) can be written as a product of the form factor and the gluon distribution

$$B(t, x_P, \bar{Q}^2) = F_B(t) \left( x_P G(x_P, \bar{Q}^2) \right). \quad (47)$$

As the pomeron-proton vertex might be similar to the photon-proton coupling [8], we shall use a simple approximation

$$F_B(t) \sim F_p^{em}(t), \quad (48)$$

where  $F_p^{em}(t)$  is the standard form for the electromagnetic form factor of the proton

$$F_p^{em}(t) = \frac{(4m_p^2 + 2.8|t|)}{(4m_p^2 + |t|)(1 + |t|/0.7\text{GeV}^2)^2}. \quad (49)$$

It has been shown in (43, 45) that the leading contribution to the  $T^+$  and  $T^-$  amplitudes is determined by the same loop integral  $I$ . For simplicity we shall suppose that the  $A$  amplitude can be parameterized in the form similar to (47)

$$A(t, x_P, \bar{Q}^2) = \alpha F_A(t) \left( x_P G(x_P, \bar{Q}^2) \right). \quad (50)$$

As previously, the new form factor  $F_A(t)$  describes the  $t$ -dependence of the two-gluon coupling with the proton for the  $A$  function. We shall use for simplicity  $F_A = F_p^{em}$ . For the approximation (50), the ratio of the  $A$  and  $B$  amplitudes is independent of  $x$ . The  $\alpha = A/B$  ratio determines through the  $x$  dependences of the functions ( $x \sim 1/s$  in this case) the energy behaviour of the spin asymmetries in exclusive reactions at high energies and fixed momentum transfer. Thus, (47) and (50) result in energy independence of spin asymmetries which is in agreement with their weak energy dependence obtained in the models [21, 22, 24]. To study the  $\alpha$  sensitivity of the cross section we shall use

in our estimations the value of  $|\alpha| \leq 0.1\text{GeV}^{-1}$ . It has been mentioned above that this magnitude is consistent with the model estimations [21, 22].

The energy dependence of the cross sections is determined by the pomeron contribution to the gluon distribution function at small  $x$

$$(x_P G(x_P, \bar{Q}^2)) \sim \left( \frac{sy}{m_J^2 + Q^2 + |t|} \right)^{(\alpha_p(t)-1)}. \quad (51)$$

Here  $\alpha_p(t)$  is a pomeron trajectory. The linear approximation of the pomeron trajectory (1) is used. The parameters of this trajectory was determined from the fit of the diffractive  $J/\Psi$  production by ZEUS [29]

$$\alpha_p(t) = 1 + (0.175 \pm 0.026) + (0.015 \pm 0.065)t. \quad (52)$$

In our model estimations of the polarized cross section of the diffractive  $J/\Psi$  production we shall use the values  $\epsilon = 0.15$  and  $\alpha' = 0$  which are not far from (52). The typical scale of the reaction is determined by (41). For not large  $Q^2$  and  $|t|$  the value of  $\bar{Q}^2$  is about 2.5-3.0  $\text{GeV}^2$ . In this region we can work with fixed  $\alpha_s \sim 0.3$ . An effective gluon mass in (38) is chosen to be equal to 0.3  $\text{GeV}^2$ . The cross section depends on this parameter weakly. The value of  $\Gamma_{e^+e^-}^J = 5.26\text{keV}$  is used.

We shall integrate the cross sections (42) over  $Q^2$  and  $y$  to get the differential cross sections of the  $J/\Psi$  production

$$\frac{d\sigma^\pm}{dt} = \int_{y_{min}}^{y_{max}} dy \int_{Q_{min}^2}^{Q_{max}^2} dQ^2 \frac{d\sigma^\pm}{dQ^2 dy dt}, \quad (53)$$

where

$$Q_{min}^2 = m_e^2 \frac{y^2}{1-y} \quad \text{and} \quad Q_{max}^2 = 4\text{GeV}^2.$$

For the HERA energy  $\sqrt{s} = 300\text{GeV}$  we integrate over the energy in the photon-proton system  $30\text{ GeV} < W_{\gamma p} < 150\text{ GeV}$  ( $W_{\gamma p} \sim \sqrt{ys}$ ) which is equivalent to the range  $0.01 < y < 0.25$ . Our results, shown in Fig. 2, describe experimental data [3] at small  $|t|$  and lie below them a little for the momentum transfer larger than 1  $\text{GeV}^2$ . This may indicate that the simple approximation of the form factor (48) used here is not good for  $|t| > 1\text{GeV}^2$ . Our estimation for the HERMES energy  $s = 50\text{GeV}^2$  is performed for integration over  $0.3 < y < 0.7$ . The predicted cross sections are shown in Fig. 3. It is seen from Figs. 2 and 3 that the spin-average cross sections are sensitive to  $\alpha$  but the shape of all curves are the same. Thus, it is difficult to extract information about the spin-dependent part of the pomeron coupling from the spin-average cross section of the diffractive vector meson production.

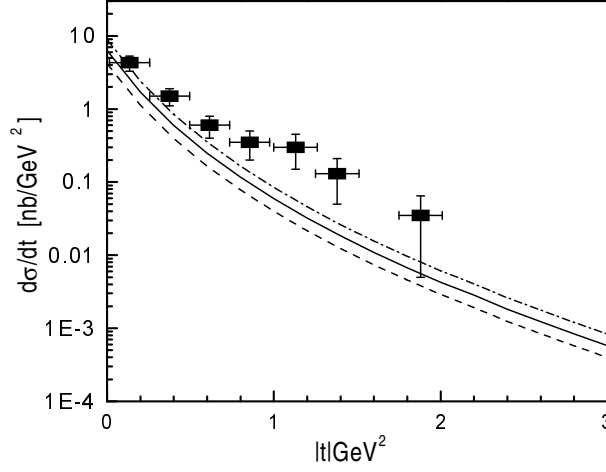


Figure 2: The differential cross section of the  $J/\Psi$  production at HERA energy: solid line -for  $\alpha = 0$ ; dot-dashed line -for  $\alpha = 0.1\text{GeV}^{-1}$ ; dashed line -for  $\alpha = -0.1\text{GeV}^{-1}$ . Data are from [3].

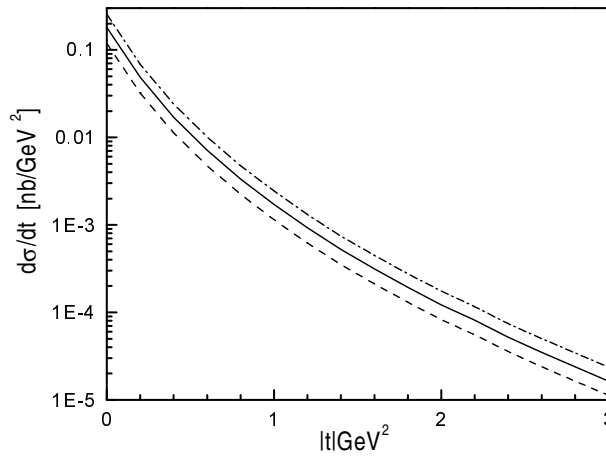


Figure 3: The differential cross section of the  $J/\Psi$  production at HERMES energy: solid line -for  $\alpha = 0$ ; dot-dashed line -for  $\alpha = 0.1\text{GeV}^{-1}$ ; dashed line -for  $\alpha = -0.1\text{GeV}^{-1}$ .

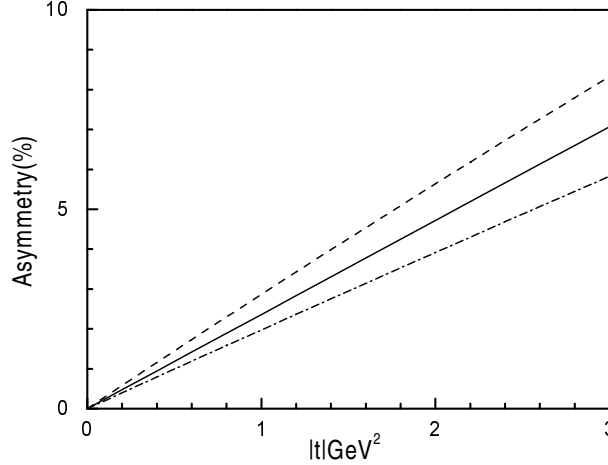


Figure 4: The predicted  $A_{ll}$  asymmetry of the  $J/\Psi$  production at HERMES: solid line -for  $\alpha = 0$ ; dot-dashed line -for  $\alpha = 0.1\text{GeV}^{-1}$ ; dashed line -for  $\alpha = -0.1\text{GeV}^{-1}$ .

Using the same formulae we calculate the cross section  $\sigma(-)$ . This gives us a possibility to estimate the longitudinal double-spin  $A_{ll}$  asymmetry of the  $J/\Psi$  production at high energies. As has been found, the asymmetry vanishes as  $1/s$  and for the HERA energy range the expected value of  $A_{ll}$  will be negligible. The predicted asymmetry for HERMES as a function  $\alpha$  is shown in Fig. 4. The important property of  $A_{ll}$  is that the asymmetry of the vector meson production is equal to zero in the forward direction. The  $A_{ll}$  asymmetry might be connected with the spin-dependent gluon distribution  $\Delta G$  only for  $|t| = 0$ . Thus,  $\Delta G$  cannot be extracted from  $A_{ll}$  in agreement with the results of [17].

To understand the  $\alpha$  dependence of the asymmetry we shall use the approximated expression of the integral (53). The functions  $d\sigma^\pm/(dQ^2 dy dt)$  decrease very rapidly with growing  $Q^2$ . Calculating only the leading Log terms of the integral over  $Q^2$  in (53) we can write the double spin asymmetry  $A_{ll}$  in a simple form

$$A_{ll} \sim \frac{|t|}{s} \frac{(2 - \bar{y})(1 + 2m\alpha)}{(2 - 2\bar{y} + \bar{y}^2)[(1 + 2m\alpha)^2 + \alpha^2|t|]}, \quad (54)$$

where  $\bar{y}$  is some average value in the integration region. We find, that the asymmetry is non zero for  $\alpha = 0$

$$A_{ll}^0 = A_{ll}(\alpha = 0) \simeq \frac{|t|}{s} \frac{(2 - \bar{y})}{(2 - 2\bar{y} + \bar{y}^2)}. \quad (55)$$

This term in asymmetry is determined by the  $\Delta A_{full}$  contribution to  $W(-)$  (see (27)). Both spin-average and spin-dependent cross sections (43, 45) are proportional to  $|B|^2 \sim (x_P G(x_P, \bar{Q}^2))$  for  $\alpha = 0$ , which results in the independence of  $A_{ll}^0$  on the gluon distribution. We see from Fig. 4 that the  $A_{ll}^0$  part of the asymmetry dominates. The value of the asymmetry for  $\alpha \neq 0$  is determined by the spin-dependent part of the pomeron coupling.



However, the sensitivity of the asymmetry to  $\alpha$  is not very strong. Thus, it will not be so easy to study the spin structure of the pomeron coupling with the proton from the  $A_{ll}$  asymmetry of the diffractive  $J/\Psi$  production.

## 6 Conclusion

In the present paper, the polarized cross section of the diffractive  $J/\Psi$  leptonproduction at high energies has been studied. The relevant cross section can be determined in terms of the leptonic and hadronic tensors; and the squared amplitude of the vector meson production, through the photon-pomeron fusion. The amplitude of the  $\gamma\mathbb{P} \rightarrow J/\Psi$  transition is described by the simple non-relativistic wave function. This approximation is efficient, at least for heavy meson production. The introduced hadronic tensor is expressed in terms of the pomeron-proton coupling structure which has the helicity flip part. As a result, connection of the spin-dependent cross section in the diffractive  $J/\Psi$  production with the pomeron coupling has been found. We predict the not small value of the  $A_{ll}$  asymmetry of the diffractive vector meson production at the HERMES energy. However, the asymmetry decreases as  $1/s$  with growing energy and at the HERA energy it will be extremely small. It has been found that the  $A_{ll}$  asymmetry vanishes at  $t = 0$ . Thus, it is impossible to extract the polarized gluon distribution  $\Delta G$  from the asymmetry. The predicted asymmetry is independent of the mass of a produced meson. We can expect a similar value of the asymmetry in the polarized diffractive  $\phi$ -meson leptonproduction.

The longitudinal double spin asymmetry of the vector meson production for nonzero momentum transfer has been found to be dependent on the  $A$  term of the pomeron coupling which produces helicity-flip effects. Note that this spin-dependent part of the coupling  $A$  is parametrized here by the gluon structure function of the proton  $G$ , for simplicity. Generally, the function  $A$  should be determined by the polarized gluon distribution and the ratio  $\alpha$  might depend on  $x_P$  and  $t$ . However, this conformity is not known quite well now. To find the explicit connection of  $A$  with spin-dependent gluon distribution in QCD, additional study is necessary. Our results show the essential role of the "full block asymmetry" in  $A_{ll}$ . This contribution does not depend on the gluon distribution and has a kinematic character to all appearance. The information about the spin-structure of the pomeron coupling can generally be extracted from the  $A_{ll}$  asymmetry of the vector meson production for  $|t| \neq 0$ . Such investigations can be carried out in future polarized experiments at CERN and DESY. However, the sensitivity of asymmetry to the  $\alpha$ -ratio is quite weak. Thus, the diffractive vector meson production might not be

a good tool to study the polarized gluon distributions of the proton and spin structure of the pomeron.

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